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ATOMS, MOLECULES, OPTICS

Light-Induced Cross Transport Phenomena in a Single-Component Gas

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Abstract—The cross transport processes that occur in a single-component gas in a capillary and are caused by resonance laser radiation and pressure and temperature gradients are studied. An expression for entropy production is derived using a system of kinetic Boltzmann equations in a linear approximation. The kinetic coefficients that determine the transport processes are shown to satisfy the Onsager reciprocal relations at any Knudsen numbers and any character of the elastic interaction of gas particles with the capillary surface. The light-induced baro- and thermoeffects that take place in a closed heat-insulated system in the field of resonance laser radiation are considered. Analytical expressions are obtained for the Onsager coefficients in an almost free-molecular regime. The light-induced pressure and temperature gradients that appear in a closed heat-insulated capillary under typical experimental conditions are numerically estimated.

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1. INTRODUCTION

Light-induced drift (LID) [1] and heat transfer (LIHT) [2, 3] are known to appear in a gas during the selective absorption of optical radiation by molecules or atoms at certain velocities. Experimental recording of the LID rate is a complex technical problem. It is much easier to measure the pressure drop at the ends of a closed capillary that appears due to gas drift along this capillary, the so-called light-induced baroeffect [4]. In a certain sense, this phenomenon is analogous to the thermomolecular pressure difference [5]. The difference consists only in the fact that LID plays the role of thermal creep.

As for the possibility of an experimental study of a light-induced heat flux, it could be measured by detecting the temperature drop at the ends of a heat-insulated capillary, the so-called light-induced thermoeffect. Imagine a capillary with its ends closed by glasses that are transparent for light. The gas in the space between the glasses is pumped to an ultrahigh vacuum, and the lateral surface is heat-insulated. In this case, the gas in the capillary can be considered as a heat-insulated system to a good approximation. If resonance electromagnetic radiation is directed along the capillary, longitudinal temperature redistribution takes place due to light-induced heat transfer; that is, a certain stationary temperature gradient occurs in the gas. What is its magnitude and can it be detected? This question can only be answered when solving the corresponding problem using the methods of the kinetic theory of gases.

If the gas temperature were uniform, a stationary pressure gradient along the capillary would be set provided that LID is compensated by the Poiseuille flow. If a uniform gas pressure was maintained along the capillary tube, a stationary temperature gradient would occur provided LIHT is compensated by the ordinary heat flux proportional to the temperature gradient. Note that LIHT is caused by the difference between the accommodation coefficients and the collision cross sections of excited and unexcited particles, as for the transfer of the translational energy of particles [3], and also includes the excitation energy flux. In a closed heat-insulated capillary, stationary pressure and temperature gradients form as a result of the interaction of LID, the Poiseuille flow, thermal creep, LIHT, ordinary heat flux, and isothermal heat transfer proportional to the pressure gradient (mechanocaloric flow) [5]. Therefore, it is interesting to develop a theory to reveal the relation between the characteristics of the interaction of resonance laser radiation and the stationary pressure and temperature gradients.

Sharipov [6] considered the transport phenomena in a mixture of optically active and buffer gases in a channel in the field of laser radiation and added pressure, temperature, and concentration gradients as sources of nonequilibrium. Based on the fundamental properties of kinetic equations and the law of gas–surface interaction, he proved the Onsager reciprocal relations for the kinetic coefficients that characterize the heat-and-mass transfer processes in the gas through the channel. However, this proof is thought to be based on insufficiently grounded assumptions; in particular, Sharipov assumed the possibility of radiative decay of the ground states in atoms and spontaneous transition from the ground state into an excited one. Moreover, the assumptions regarding the relax-

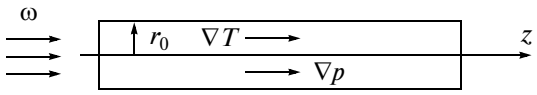

Fig. 1. Geometry of the problem.

ation of the populations of the ground and excited levels eventually led to the conclusion that the kinetic coefficients are independent of the frequency of radiative decay of the excited level. This conclusion is in conflict with the result [7] obtained by a direct numerical solution of kinetic equations.

The purpose of this work is to develop a physicomathematical model for the light-induced cross phenomena that take place in the gas located in a capillary and having a nonuniform pressure and temperature. To accomplish this purpose, we solve the following problems.

(1) Based on linearized Boltzmann equations for excited and unexcited gas particles, we first obtain an expression for entropy production in the gas–capillary system using general boundary conditions for distribution functions and then choose generalized flows and forces.

(2) We derive expressions for the kinetic coefficients through the distribution functions of excited and unexcited particles and prove the symmetry of the cross kinetic coefficients for any Knudsen (Kn) numbers and any character of the elastic interaction of gas particles with the wall.

(3) We consider the stationary state of the closed gas–capillary system at a fixed generalized force related to radiation. Based on the Prigogine theorem about stationary states with the minimum entropy production [5], we obtain expressions for the pressure and temperature gradients that form along the capillary in terms of the kinetic coefficients.

(4) We solve kinetic equations with the boundary conditions of specular–diffusion reflection of particles from the capillary wall for the case of an almost free-molecular regime. With this solution, we derive analytical expressions for the kinetic coefficients and the stationary gas pressure and temperature gradients.

2. FORMULATION OF THE PROBLEM

We consider a single-component gas in a long capillary of radius r_0. Resonance optical radiation propagates along the capillary axis in the form of a traveling monochromatic wave, and longitudinal gas pressure and temperature gradients are present.

We choose cylindrical coordinate system (r, φ, z) so that polar axis z is directed along the capillary axis (Fig. 1).

We use the approximation according to which stimulated resonance transitions only between two "working" levels, i.e., ground (n) and excited (m) ones, are possible in multilevel particles. Transition frequency ω_{nm} is close to radiation frequency ω. Because of the Doppler effect, radiation is only absorbed by the particles whose longitudinal component of velocity vector $\mathbf{v}$ differs weakly from resonance value $kv_{z,\,\mathrm{res}} = \Omega = \omega - \omega_{mn}$, where k is the wavenumber and Ω is the radiation frequency detuning ($|\Omega| \ll \omega, \omega_{mn}$). The particles absorbed radiation change their transport properties, in particular, the collision cross section. Thus, the gas can be considered as a mixture of particles having the same mass m but different interaction cross sections. Simultaneously with stimulated transitions, the radiative decay of excited level m occurs at frequency Γ_m. A continuous particle exchange between the components takes place. It is known [1] that partial light-induced mass and heat fluxes appear in such a model mixture and that a drift of the gas as a whole and heat transfer appear under certain conditions.

The distribution functions of excited (f_m) and unexcited (f_n) particles satisfy the kinetic Boltzmann equations [8]

$$\begin{aligned}
&\frac{\partial f_n}{\partial t} + \mathbf{v}_\perp \frac{\partial f_n}{\partial \mathbf{r}} + v_z \frac{\partial f_n}{\partial z} \\
&\quad = \frac{\Gamma_m \chi(\mathbf{v})}{2}(f_m - f_n) + \Gamma_m f_m + S_n, \\
&\frac{\partial f_m}{\partial t} + \mathbf{v}_\perp \frac{\partial f_m}{\partial \mathbf{r}} + v_z \frac{\partial f_m}{\partial z} \\
&\quad = \frac{\Gamma_m \chi(\mathbf{v})}{2}(f_n - f_m) - \Gamma_m f_m + S_m,
\end{aligned} \tag{1}$$

$$\chi(\mathbf{v}) = \frac{4|G_{mn}|^2 \Gamma}{\Gamma_m[\Gamma^2 + (\Omega - \mathbf{k}\cdot\mathbf{v})^2]}, \quad G_{mn} = \frac{Ed_{mn}}{2\hbar},$$

$$S_n = S_{nn} + S_{nm}, \quad S_m = S_{mn} + S_{mm}, \quad v_\perp^2 = v_r^2 + v_\varphi^2.$$

Here, $\mathbf{v}_\perp$ is the component of the particle velocity vector in the cross section of the capillary, Γ is the homogeneous absorption line half-width, E is the electric field amplitude of the light wave, d_{mn} is the dipole moment of the m–n transition, $\hbar$ is Plank's constant, $\mathbf{k}$ is the wavevector, and S_{ij} are the integrals of collisions between the particles of the ith and jth kinds. Rabi frequency G_{mn} determines the rate of the stimulated transitions of particles between levels per unit time. Saturation parameter $\chi(\mathbf{v})$ characterizes the probability of radiation absorption by particles moving at velocity $\mathbf{v}$ per unit time.

The collisions of atoms with the capillary surface are assumed to be elastic. Then, the boundary conditions for the distribution functions of excited and unexcited particles have the form [9]

$$|\mathbf{v}\cdot\mathbf{n}| f_i^+(\mathbf{v})$$

$$= \int\limits_{(\mathbf{v}'\mathbf{n})<0} R(\mathbf{v}' \longrightarrow \mathbf{v}) f_i^-(\mathbf{v}') |\mathbf{v}' \cdot \mathbf{n}| d\mathbf{v}', \qquad (2)$$

$$\mathbf{v} \cdot \mathbf{n} > 0; \quad i = m, n,$$

where $\mathbf{n}$ is the internal normal to the capillary surface; f_i^+ and f_i^- are the distribution functions of reflected particles of the ith kind and those incident on the wall, respectively; and $R(\mathbf{v}' \longrightarrow \mathbf{v})$ is the scattering function (i.e., the probability density of the fact that velocity $\mathbf{v}'$ of a particle incident on the wall is $\mathbf{v}$ after reflection).

In the weak-field approximation ($\|\chi(\mathbf{v})\| \ll 1$) and at small gradients of thermodynamic parameters, the distribution functions of excited and unexcited particles differ insignificantly from the local equilibrium Maxwell–Boltzmann distributions

$$f_i = f_{i0}[1 + h_i(\mathbf{r}, \mathbf{v})], \quad \|h_i\| \ll 1,$$

$$f_{i0} = n_i(z)\left(\frac{m}{2\pi k_\mathrm{B} T(z)}\right)^{3/2} \exp\left(-\frac{m v^2}{2k_\mathrm{B} T(z)}\right), \qquad (3)$$

$$n_i(z) = n(z)\frac{\exp(-E_i/k_\mathrm{B}T(z))}{\sum\limits_l \exp(-E_l/k_\mathrm{B}T(z))}.$$

Here, $T(z)$, $n(z)$, and $n_i(z)$ are the local value of the gas temperature, the total particle number density, and the population of the ith level, respectively; k_B is the Boltzmann constant; and E_i is the internal energy of the ith state of a particle. Note that the dependence of perturbation function $h_i(\mathbf{r}, \mathbf{v})$ on coordinate z may be neglected only in the case of a sufficiently long capillary (the length is much larger than the radius) and an optically thin medium.

The authors of [8] noted that, for a homogeneous broadening of a absorption line ($\Gamma \gg k\bar{v}$, $\bar{v} = (2k_\mathrm{B}T_0/m)^{1/2}$, T_0 is the equilibrium temperature, $k\bar{v}$ is the Doppler broadening), kinetic equations can be linearized in both local and integral senses. Velocity-space-averaged values are of interest in this work, and not the distribution function itself. Therefore, the integrated conditions of linearization are important in this case. There exists a criterion that not always meets a linearization condition in a local sense but meets it in an integral sense. For an inhomogeneous broadening of an absorption line ($\Gamma \ll k\bar{v}$), this criterion consists in the requirement that the relative difference between the nonequilibrium and equilibrium (Boltzmann) populations of the excited level is small.

With allowance for these assumptions, Eqs. (1) are linearized with respect to perturbation function h_i and are written as

$$\frac{\partial h_i}{\partial t} + \mathbf{v}_\perp \frac{\partial h_i}{\partial \mathbf{r}} = \sum_{j=n,m} L_{ij}(h_i) + R_i h_m + \frac{\bar{v}}{r_0} g_i, \qquad (4)$$

where

$$g_i = -c_z \nu - c_z\left(c^2 - \frac{5}{2} + \varepsilon_i - \bar{\varepsilon}\right)\tau + \frac{\Gamma_m r_0}{2\bar{v}}\chi(v_z)\left(\frac{n_{j0}}{n_{i0}} - 1\right), \quad i,j = n,m; \quad i \neq j, \qquad (5)$$

$$\nu = \frac{r_0}{p}\frac{\partial p}{\partial z}, \quad \tau = \frac{r_0}{T}\frac{\partial T}{\partial z},$$

$$c_i = \frac{v_i}{\bar{v}}, \quad c^2 = c_z^2 + c_\perp^2, \quad \varepsilon_i = \frac{E_i}{kT_0},$$

$$R_i = \begin{cases} \Gamma_m\left(\dfrac{n_{m0}}{n_{n0}}\right), & i = n, \\ -\Gamma_m, & i = m, \end{cases}$$

$$\bar{\varepsilon} = \frac{\sum\limits_l \varepsilon_l \exp(-\varepsilon_l)}{\sum\limits_l \exp(-\varepsilon_l)} = \frac{\bar{E}}{kT_0}.$$

Here, $L_{ij}(h_i)$ is the linearized collision integral for particles of the ith and jth kinds; n_{n0} and n_{m0} are the equilibrium populations of the ground and excited states at equilibrium gas temperature T_0, respectively; and $\bar{E}$ is the average internal energy of particles in the equilibrium state.

3. ENTROPY PRODUCTION

We mentally separate gas volume V that is present in part of the capillary of unit length and determine its entropy as usually [10],

$$S_e = -k_\mathrm{B}\sum_i \int dV \int f_i \ln f_i d\mathbf{v}. \qquad (6)$$

The rate of change of entropy is

$$\frac{\partial S_e}{\partial t} = -k_\mathrm{B}\sum_i \int dV \int (\ln f_i + 1)\frac{\partial f_i}{\partial t} d\mathbf{v}. \qquad (7)$$

We substitute the expression for $\partial f_i/\partial t$ from Eqs. (1) into this equation; use the linearization of distribution functions (3); take into account the equations of continuity, motion, and the conservation of the total inter-

nal energy of the gas volume under study; assume $\ln(1 + h_i) \approx h_i - h_i^2/2$; and obtain

$$\frac{\partial S_e}{\partial t} = \sigma - k_B \frac{\bar{V}}{r_0} \sum_i \int_S dS \int f_{i0} h_i g_i d\mathbf{v} - \int_V \frac{1}{T} \mathrm{div}\mathbf{q}\, dV, \quad (8)$$

where S is the cross-sectional area of the capillary and $\mathbf{q}$ is the heat flux density vector in the gas. The first term in the right-hand side of Eq. (8) (σ) determines the rate of increase of entropy in the separated gas volume, and the second term is the rate of change of entropy in the gas due to its interaction with the environment. To reveal the meaning of the third term, we transform it. To this end, we take into account that the condition of continuity of normal heat flux $q_n = q_{wn}$, where $q_n = \mathbf{q} \cdot \mathbf{n}$, $q_{wn} = \mathbf{q}_w \cdot \mathbf{n}$ are the normal components of the heat flux density vectors in the gas and the capillary wall, respectively, is met at any surface point in the stationary state, $\mathbf{n}$ is the inward normal to the wall. Let $T_w(z)$ be the local temperature of the inner capillary surface. Then, using the divergence theorem and the law of conservation of the capillary internal energy, we have

$$\int_V \frac{1}{T} \mathrm{div}\mathbf{q}\, dV = \int_{S_w} \frac{q_n}{T_w} dS = \int_{S_w} \frac{q_{wn}}{T_w} dS, \quad (9)$$

where S_w is the lateral area of the capillary. Thus, the third term in the right-hand side of Eq. (8) takes into account the entropy production inside the capillary walls.

The entropy production in the gas is written in the form $\sigma = \sigma_{\mathrm{coll}} + \sigma_R + \sigma_w$; that is, it is determined by the following three processes:

molecular collisions,

$$\sigma_{\mathrm{coll}} = -k_B \sum_{(i,j)=n,m} \int_V dV \int f_{i0} h_i L_{ij}(h_i) d\mathbf{v}; \quad (10)$$

radiative decay of the excited level,

$$\sigma_R = k_B \frac{\Gamma_m}{2} \int_V dV \int f_{m0} (h_m - h_n)^2 d\mathbf{v}; \quad (11)$$

interaction of gas particles with the capillary surface,

$$\sigma_w = k_B \sum_{i=n,m} \int_V dV \int f_{i0} h_i \mathbf{v}_\perp \frac{\partial h_i}{\partial \mathbf{r}} d\mathbf{v} = k_B \sum_{i=n,m} \int_\Sigma d\Sigma \int (\mathbf{n} \cdot \mathbf{v}_\perp) f_{i0} \frac{h_i^2}{2} d\mathbf{v}. \quad (12)$$

Here, $\mathbf{n}$ is the internal normal to the surface Σ that covers gas volume V under study and includes the cross-sectional area S and the surface area S_w of the capillary. In the last equality of Eq. (12), we used the Gauss-Ostrogradski formula. It follows from the symmetry properties of the collision integral that $\sigma_{\mathrm{coll}} \geq 0$ [10]. According to Eq. (11), we have $\sigma_R \geq 0$. Using the reciprocal relation for the scattering probability density $R(\mathbf{v}' \longrightarrow \mathbf{v})$ and the Jensen inequality for convex downward functions, we have $\sigma_w \geq 0$ [9]. Thus, the entropy production in the gas is not negative, $\sigma \geq 0$.

In the stationary state, the entropy of the gas volume under study does change in time. Then it follows from Eq. (8) that

$$\sigma = k_B \frac{\bar{V}}{r_0} \sum_{i=m,n} \int_S dS \int f_{i0} h_i g_i d\mathbf{v} + \int_{S_w} \frac{q_{wn}}{T_w} dS. \quad (13)$$

The contribution of the second term to the entropy production depends on both the thermal conductivity of the capillary material and the ratio of the capillary radius r_0 to its length L [11, 12]. If $L \gg r_0$, the contribution of the second term in the right-hand side of Eq. (13) becomes negligibly small at any thermal conductivity of the wall. In the limit $L/r_0 \longrightarrow \infty$, the state of local equilibrium, where the normal component of the heat flux density vector is zero and the wall temperature is equal to the gas temperature, forms in the cross section of the capillary in the stationary case. Since a long capillary ($L \gg r_0$) is considered, we neglect the contribution of the thermal conductivity of the capillary wall to the entropy production; that is, the second term in the right-hand side of Eq. (13) is omitted.

Following the thermodynamics of irreversible processes [5], we represent the entropy production as the sum of the products of generalized flows J_k into conjugate generalized forces X_k,

$$\sigma = J_p X_p + J_T X_T + J_S X_S. \quad (14)$$

The generalized forces are chosen as follows:

$$X_p = -k_B \nu, \quad X_T = -\frac{\tau}{T_0}, \quad X_S = \chi_0 = \int_{-\infty}^{\infty} \chi(c_z) e^{-c_z^2} dc_z. \quad (15)$$

The first two forces were used in [13] to describe the nonisothermal gas motion in a channel. The choice of force X_S related to resonance optical radiation is convenient, since quantity χ_0 has the meaning of the probability of light absorption by gas particles per unit time and is measured in spectroscopic experiments. Note that force χ_0 is proportional to radiation intensity I, is independent of the sign of detuning Ω, and becomes maximal under the exact resonance condition ($\Omega = 0$).

As follows from Eqs. (13) and (14) with allowance for the form of source function g_i in Eq. (5), general-

ized flows J_k conjugate to forces (15) are described by the expressions

$$J_p = (\varphi_p, h_i), \quad J_T = k_B T_0(\varphi_T, h_i), \quad J_S = k_B(\varphi_S, h_i), \tag{16}$$

where

$$\varphi_p = \frac{c_z \bar{V}}{r_0}, \quad \varphi_T = \frac{c_z \bar{V}}{r_0}\left(c^2 - \frac{5}{2} + \varepsilon_i - \bar{\varepsilon}\right),$$

$$\varphi_S = \Gamma_m \frac{\chi(c_z)}{2\chi_0}\left(\frac{n_{j0}}{n_{i0}} - 1\right), \tag{17}$$

$$(\varphi_l, h_i) = \sum_{i=n,m} \int dS \int \varphi_l f_{i0} h_i d\mathbf{v},$$

$$l = p, T, S, \quad i, j = n, m; \quad i \neq j.$$

The flows determined in Eq. (16) have the following obvious physical meaning: J_p and J_T are the gas and heat flow densities averaged over the cross section of the capillary, respectively, and J_S is the entropy flow density conjugate to force X_S.

4. KINETIC COEFFICIENTS

In a linear approximation, the perturbation function can be represented as a linear combination of the generalized forces,

$$h_i = h_i^p \nu + h_i^T \tau + h_i^S \chi_0. \tag{18}$$

Substituting Eq. (18) into Eqs. (16), for the flows we obtain

$$J_l = \sum_k \Lambda_{lk} X_k, \quad (k, l) = p, T, S, \tag{19}$$

where

$$\Lambda_{pp} = -\frac{1}{k_B}(\varphi_p, h_i^p), \quad \Lambda_{pT} = -T_0(\varphi_p, h_i^T),$$

$$\Lambda_{pS} = (\varphi_p, h_i^S), \quad \Lambda_{Tp} = -T_0(\varphi_T, h_i^p),$$

$$\Lambda_{TT} = -k_B T_0^2(\varphi_T, h_i^T), \quad \Lambda_{TS} = k_B T_0(\varphi_T, h_i^S),$$

$$\Lambda_{Sp} = -(\varphi_S, h_i^p), \quad \Lambda_{ST} = -k_B T_0(\varphi_S, h_i^T), \quad \Lambda_{SS} = k_B(\varphi_S, h_i^S). \tag{20}$$

Here, Λ_{lk} is the kinetic coefficient that determines the contribution of the component proportional to generalized force X_k to flow J_l. In so-called forward processes, force X_k induces the component of conjugate flow J_k that is characterized by coefficient Λ_{kk}. Forward processes are represented by Poiseuille flow Λ_{pp} and heat flux Λ_{TT}, which are caused by the pressure and temperature gradient, respectively, and force X_S causes the entropy flow along the capillary due to stimulated transitions Λ_{SS}. In addition, forces X_k stimulate nonconjugated flows J_i ($i \neq k$), which are characterized by kinetic coefficients Λ_{ik}; that is, they are cross flows. In particular, the temperature gradient stimulates gas motion along the capillary (thermal creep Λ_{pT}) and entropy flux (thermoentropy flux Λ_{ST}), and the pressure gradient induces heat flux (mechanocaloric flux Λ_{Tp}) [5] and entropy flux (baroentropy flux Λ_{Sp}). The action of laser radiation causes gas drift Λ_{pS} and heat flux Λ_{TS}, which are cross flows with respect to the baro- (Λ_{Sp}) and thermoentropy (Λ_{ST}) fluxes, respectively.

In the thermodynamics of nonequilibrium processes, the following Onsager reciprocal relations are postulated for cross coefficients:

$$\Lambda_{lk} = \Lambda_{kl}. \tag{21}$$

To prove the symmetry of Onsager matrix (21), we use the obvious principle of a balance of the number of particles, according to which the decrease in the population of excited level m for particles with velocity $\mathbf{v}$ that is caused by radiative decay is equal to (by an order of magnitude) the increase in the population of ground state n for the particles moving at the same velocity. By substituting Eq. (18) into Eq. (4) and choosing terms at the same generalized forces, we obtain a system of kinetic equations for each perturbation function h_i^k ($i = n, m$; $k = p, T, S$). From an analysis of these equations (which differ only in free terms from each other), we have

$$h_i^p\left(c^2 - \frac{5}{2} + \varepsilon_i - \bar{\varepsilon}\right) = h_i^T, \quad h_i^p \varphi_S = -\varphi_p h_i^S, \quad h_i^T \varphi_S = -\varphi_T h_i^S. \tag{22}$$

With allowance for Eq. (22), Eqs. (20) for cross kinetic coefficients Λ_{lk} satisfy Onsager reciprocal relations (21). Note that this result holds true for any Kn numbers and any character of the elastic interaction of gas particles with the capillary surface.

5. LIGHT-INDUCED BARO- AND THERMOEFFECTS

We consider the system that consists of a single-component gas in a closed capillary with thermally insulated walls. In this case, the ends of the capillary can transmit laser radiation.

Let force X_S, which is related to resonance laser radiation, be fixed. As a result of light-induced gas drift and heat flux, longitudinal pressure and temperature gradients appear in the system. The transport phenomena stimulated by them tend to suppress the light-induced processes that caused them. As a result of increasing competition of nonequilibrium processes, a certain stationary state of dynamic equilibrium forms eventually in the system. According to the Prigogine principle [5], this state is characterized by the minimum entropy production, where flows J_p and

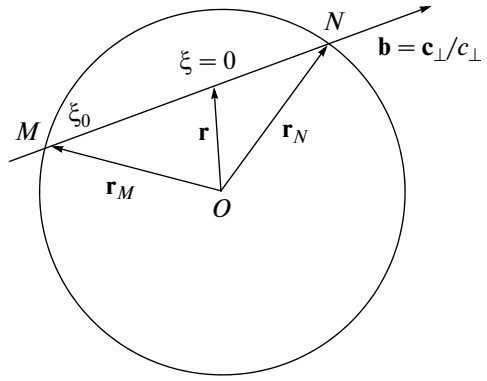

Fig. 2. Schematic diagram for the integration of the kinetic equation along a particle velocity vector.

J_T conjugate to unfixed forces X_p and X_T should be zero,

$$\left(\frac{\partial\sigma}{\partial X_p}\right)_{X_S} = J_p$$
$$= \Lambda_{pp}X_p + \Lambda_{pT}X_T + \Lambda_{ps}X_s = 0,$$
$$\left(\frac{\partial\sigma}{\partial X_T}\right)_{X_S} = J_T \quad (23)$$
$$= \Lambda_{Tp}X_p + \Lambda_{TT}X_T + \Lambda_{Ts}X_s = 0.$$

From these equations, we obtain the following expressions for the pressure and temperature gradients:

$$X_p = -k_\mathrm{B}\nu = X_S(\Lambda_{pT}\Lambda_{TS} - \Lambda_{pS}\Lambda_{TT})/\Delta_0,$$
$$X_T = -\frac{\tau}{T_0} = X_S(\Lambda_{Tp}\Lambda_{pS} - \Lambda_{pp}\Lambda_{TS})/\Delta_0, \quad (24)$$

where

$$\Delta_0 = \Lambda_{pp}\Lambda_{TT} - \Lambda_{Tp}\Lambda_{pT}.$$

It is seen that the longitudinal gas pressure redistribution in the capillary is caused by LID and thermal creep and depends on the heat fluxes forming a temperature nonuniformity, which, in turn, causes thermal creep. The temperature gradient is induced by LIHT and the mechanocaloric flow and depends on the Poiseuille flow and LID, which form the pressure gradient and, hence, the mechanocaloric flow.

Thus, resonance laser radiation can serve as a piston compressing the gas and a specific heat engine that can redistribute the heat energy in the gas.

Coefficients Λ_{kl} determined by Eqs. (20) can be calculated from the solution to kinetic equations (4) using a certain model of boundary conditions (2).

6. APPROXIMATION OF ALMOST FREE-MOLECULAR REGIME

An analytical calculation of the kinetic coefficients can be rather simply performed in the case of an almost free-molecular regime (Kn $\gg$ 1). We use the successive collisions approximation [14]; according to it, a particle leaves an element of the phase volume after the first collision and no reverse collision, which results in the motion of a new particle into this volume, takes place. Thus, the reverse collision integrals in Eq. (4) disappear and the linearized collision integrals can be approximated as follows:

$$L_{ij}(h_i) = -\gamma_{ij}h_i, \quad (25)$$

where γ_{ij} are the effective frequencies of collisions between the ith and jt particles that are independent of the molecular velocities. Note that the successive collisions method gives good results for sufficiently high Kn numbers [14].

As boundary conditions for distribution functions, we use the Maxwell model of specular–diffusion reflection. We assume that, during an elastic collision with the capillary surface, the fraction of ith atoms (α_i) is scattered in a diffusion manner with a Maxwellian velocity distribution and fraction $1 - \alpha_i$ is specularly reflected. In the general case, the accommodation coefficients of excited (α_m) and unexcited (α_n) particles are different. Then, boundary conditions (2) are written in the form

$$f_i^+(\mathbf{v}) = \alpha_i f_i^w(\mathbf{v}) + (1-\alpha_i)f_i^-(\mathbf{v} - 2(\mathbf{v}\cdot\mathbf{n})\mathbf{n}),$$
$$\mathbf{v}\cdot\mathbf{n} > 0,$$
$$f_i^w(\mathbf{v}) = n_i^w\left(\frac{m}{2\pi k_\mathrm{B}T_w}\right)^{3/2}\exp\left(-\frac{m\mathbf{v}^2}{2k_\mathrm{B}T_w}\right), \quad (26)$$
$$i = m, n;$$

and superscripts +, w, and − belong to the reflected, diffusionally scattered, and incident particles, respectively, and n_i^w is the number density of diffusionally scattered ith particles. This density corresponds to local wall temperature T_w, which is gas temperature $T(z)$ in the corresponding section of the capillary in the assumed approximation ($L \gg r_0$). With allowance for linearization (3), boundary conditions (26) at certain point M in the capillary wall for perturbation functions take the form

$$h_i^+(\mathbf{r}_M, \mathbf{v}) = (1-\alpha_i)h_i^-(\mathbf{r}_M, \mathbf{v} - 2(\mathbf{v}\cdot\mathbf{n})\mathbf{n}),$$
$$\mathbf{v}\cdot\mathbf{n} > 0, \quad i = m, n, \quad (27)$$

where $\mathbf{r}_M$ is the radius vector of point M, $|\mathbf{r}_M| = r_0$, and $\mathbf{n}$ is the internal normal at point M.

To solve kinetic equations (4) with collision integrals in form (25) with boundary conditions (27), we use integration along the transverse component of the particle velocity vector $\mathbf{b} = \mathbf{c}_\perp/c_\perp$ (Fig. 2). In this case, boundary conditions (27) relate the distribution function of the particles emitted from point M in the capillary surface toward vector $\mathbf{b}$ to the distribution function of the particles incident on the channel surface at this point with velocity direction $-\mathbf{b}$.

The symmetry condition implies the equality of the perturbation functions of the particles incident on the capillary wall at points M and N along ray $\mathbf{b}$,

$$h_i^-(\mathbf{r}_M, -\mathbf{b}) = h_i^-(\mathbf{r}_N, \mathbf{b}). \tag{28}$$

As a result of the integration of Eqs. (4) along vector $\mathbf{b}$ and using boundary conditions (27) and symmetry conditions (28), we obtain

$$h_i^{(l)}(\mathbf{r}, \mathbf{c}_\perp) = \pm\frac{r_0\varphi_l}{\bar{v}\delta_i}\Bigg\{1 - \exp\left(-\frac{\delta_i\xi_0}{c_\perp}\right) + (1-\alpha_i)$$
$$\times\left[\exp\left(-\frac{\delta_i\xi_0}{c_\perp}\right) - \exp\left(-\frac{\delta_i(\xi_0+\vartheta)}{c_\perp}\right)\right]$$
$$\times\left[1-(1-\alpha_i)\exp\left(-\frac{\delta_i\vartheta}{c_\perp}\right)\right]^{-1}\Bigg\},$$
$$l = p, T, s; \quad i = m, n,$$
$$\xi_0 = |\mathbf{r}_M - \mathbf{r}|, \quad \vartheta = |\mathbf{r}_M - \mathbf{r}_N|, \tag{29}$$
$$\delta_m = \delta(\theta + \Gamma_m^n), \quad \delta_n = \delta(1+\Gamma_m^n),$$
$$\Gamma_m^{(n)} = \frac{\Gamma_m}{\gamma_{nn}+\gamma_{nm}},$$
$$\theta = \frac{\gamma_{mm}+\gamma_{mn}}{\gamma_{nn}+\gamma_{nm}}, \quad \delta = \frac{r_0(\gamma_{nn}+\gamma_{nm})}{\bar{v}}.$$

In the right-hand side of Eq. (29), sign + is chosen at $l = s$ and sign − is chosen at $l = p, T$.

Gas particles are simulated as hard spheres with effective diameters d_i ($i = m, n$). Since the difference between the effective diameters of the excited and unexcited particles is small [15], we introduce a small parameter

$$\left|\frac{\Delta d}{d_n}\right| = \left|\frac{d_m - d_n}{d_n}\right| \ll 1. \tag{30}$$

Hence, allowing for the small population of the excited level ($n_m \ll n_n$), we have

$$\theta = \frac{\gamma_{mm}+\gamma_{mn}}{\gamma_{nn}+\gamma_{nm}} = \frac{d_{mn}^2}{d_n^2} \approx 1 + \frac{\Delta d}{d_n},$$
$$d_{mn} = \frac{1}{2}(d_m + d_n).$$

It is also known that the coefficients of particle accommodation on the capillary walls are close to unity [4, 16]. Therefore, we consider the case of almost diffusion scattering of atoms by the channel surface,

$$1 - \alpha_i \ll 1, \quad i = m, n. \tag{31}$$

If collision frequency γ_{nn} is chosen to be $\gamma_{nn} = p/\eta$ and coefficient of viscosity η is determined as $\eta = \rho\bar{v}l/\sqrt{\pi}$ (where l is the mean free path and ρ is the gas density), rarefaction parameter δ is related to the Knudsen number by the expression

$$\delta = \frac{\sqrt{\pi}}{2}\mathrm{Kn}^{-1}, \quad \mathrm{Kn} = \frac{l}{r_0}. \tag{32}$$

Using linearization of the problem in small parameters (30) and (31), we can separate the surface and collisional mechanisms of the light-induced transport phenomena [3].

As a result, we derive the following expressions for the kinetic coefficients in the case of an almost free-molecular regime ($\delta \ll 1$):

$$\Lambda_{pp} = \frac{n_0\bar{v}}{k_\mathrm{B}}\left[\frac{4}{3\sqrt{\pi}} + \frac{1}{2}G_1 + (1-\alpha_n)\right.$$
$$\left.\times\left(\frac{8}{3\sqrt{\pi}} + 3G_1\right)\right], \tag{33}$$

$$\Lambda_{pT} = \Lambda_{Tp} = -n_0\bar{v}T_0\left[\frac{2}{3\sqrt{\pi}} + \frac{1}{2}G_1 + (1-\alpha_n)\right.$$
$$\left.\times\left(\frac{4}{3\sqrt{\pi}} + 3G_1\right)\right], \tag{34}$$

$$\Lambda_{TT} = p_0\bar{v}T_0\left[\frac{3}{\sqrt{\pi}} + \frac{5}{4}G_1 + 3(1-\alpha_n)\right.$$
$$\left.\times\left(\frac{2}{\sqrt{\pi}} + \frac{5}{2}G_1\right)\right], \tag{35}$$

$$\Lambda_{pS} = \Lambda_{Sp} = \frac{n_0 r_0\Gamma_m}{2\sqrt{\pi}}\frac{\chi_1}{\chi_0}$$
$$\times\left[\left(\frac{16}{3\sqrt{\pi}} + 6G_1\right)\Delta\alpha + G_2\frac{\Delta d}{d_n}\right], \tag{36}$$

$$\Lambda_{TS} = \Lambda_{ST} = \frac{p_0 r_0\Gamma_m}{2\sqrt{\pi}\chi_0}\Bigg\{\left[-\chi_1\left(\frac{32}{3\sqrt{\pi}} + 15G_1\right)\right.$$
$$\left.+\chi_3\left(\frac{16}{3\sqrt{\pi}} + 6G_1\right)\right]\Delta\alpha \tag{37}$$
$$-\left(\frac{5}{2}\chi_1 - \chi_3\right)G_2\frac{\Delta d}{d_n} + \frac{\hbar\omega_{mn}}{k_\mathrm{B}T_0}\chi_1 G_3\Bigg\},$$

$$\Lambda_{SS} = \frac{p_0 r_0^2\Gamma_m^2}{4\sqrt{\pi}\bar{v}T_0}\frac{n_{n0}}{n_{m0}}\frac{\mu_0}{\chi_0^2}G_3, \tag{38}$$
$$p_0 = n_0 k_\mathrm{B} T_0, \quad \Delta\alpha = \alpha_n - \alpha_m,$$
$$G_1 = (1+\Gamma_m^n)\delta\ln\delta, \quad G_2 = \delta\ln\delta,$$
$$G_3 = \frac{8}{3\sqrt{\pi}} + G_1 + G_2\frac{\Delta d}{d_n}$$
$$+ (1-\alpha_m)\left(\frac{16}{3\sqrt{\pi}} + 6G_1\right),$$

$$\chi_k = \int_{-\infty}^{\infty} c_z^k \chi(c_z)\exp(-c_z^2)dc_z,$$
$$\mu_k = \int_{-\infty}^{\infty} c_z^k \chi^2(c_z)\exp(-c_z^2)dc_z. \quad (39)$$

Quantities χ_0, χ_1, χ_3, and μ_0 include the characteristics of radiation and its interaction with gas particles, depend on detuning $\Omega/(k\bar{v})$ and broadening $\Gamma/(k\bar{v})$, and can be expressed in terms of the plasma function the numerical values of which are tabulated in [17]. In the case of an inhomogeneous absorption line broadening ($\Gamma \ll k\bar{v}$, which is typical of gases), the expressions for χ_n and μ_0 in the zeroth approximation in parameter $\Gamma/(k\bar{v})$ have the form

$$\chi_n = \frac{4\pi}{k\bar{v}}\frac{|G_{mn}|^2}{\Gamma_m}\left(\frac{\Omega}{k\bar{v}}\right)^n \exp\left[-\left(\frac{\Omega}{k\bar{v}}\right)^2\right],$$
$$\mu_0 = \frac{2|G_{mn}|^2}{\Gamma_m\Gamma}\chi_0, \quad k = \frac{2\pi}{\lambda}, \quad (40)$$

where λ is the radiation wavelength.

In the case of a homogeneous broadening ($\Gamma \gg k\bar{v}$), we have

$$\chi_1 = \frac{4\sqrt{\pi}\Omega(k\bar{v})\Gamma}{\Gamma_m}\left(\frac{|G_{mn}|}{\Omega^2+\Gamma^2}\right)^2, \quad \chi_3 = \frac{3}{2}\chi_1,$$
$$\chi_0 = \frac{4\sqrt{\pi}|G_{mn}|^2\Gamma}{\Gamma_m(\Omega^2+\Gamma^2)}, \quad \mu_0 = \frac{4|G_{mn}|^2\Gamma}{\Gamma_m(\Omega^2+\Gamma^2)}\chi_0. \quad (41)$$

In Eq. (37), coefficient Λ_{TS}, which characterizes LIHT, includes terms proportional to $\Delta\alpha$, $\Delta d/d_n$, and $\hbar\omega_{mn}/k_B T_0$. The first term determines the accommodation LIHT mechanism; the second, the surface mechanism; and the third represents excitation energy transfer. The same is true of the thermoentropy flow (Λ_{ST}), which is cross to LIHT.

We now substitute Eqs. (33)–(37) derived for the kinetic coefficients into Eqs. (24). After linearization in small parameters $\Delta d/d_n$ and $1-\alpha_n$, we obtain the following expressions for the gas pressure and temperature gradients that form in the close heat-insulated system:

$$\frac{dp}{dz} = \frac{p_0\Gamma_m}{2\sqrt{\pi}\bar{v}}\gamma_p, \quad \frac{dT}{dz} = \frac{T_0\Gamma_m}{2\sqrt{\pi}\bar{v}}\gamma_T, \quad (42)$$

where dimensionless quantities γ_p and γ_T have the form

$$\gamma_p = \left[\chi_3 + \frac{5}{2}\chi_1 + \frac{75\sqrt{\pi}}{64}\left(\chi_3 + \frac{37}{50}\chi_1\right)G_1\right]\Delta\alpha$$
$$+\left[\frac{3\sqrt{\pi}}{16}(\chi_3+2\chi_1)G_2\right]\frac{\Delta d}{d_n} + \frac{\hbar\omega_{mn}}{k_B T_0}\chi_1$$
$$\times\left[\frac{1}{2} + \frac{27\sqrt{\pi}}{128}G_1 + \frac{3\sqrt{\pi}}{16}G_2\frac{\Delta d}{d_n}\right.$$
$$\left.+\left(1+\frac{75\sqrt{\pi}}{64}G_1\right)\Delta\alpha + (1-\alpha_n)\frac{51\sqrt{\pi}}{64}G_1\right], \quad (43)$$

$$\gamma_T = \left[2\left(\chi_3 - \frac{3}{2}\chi_1\right) + \frac{51\sqrt{\pi}}{32}\left(\chi_3 - \frac{67}{34}\chi_1\right)G_1\right]\Delta\alpha$$
$$+\left[\frac{3\sqrt{\pi}}{8}(\chi_3-2\chi_1)G_2\right]\frac{\Delta d}{d_n} + \frac{\hbar\omega_{mn}}{k_B T_0}\chi_1$$
$$\times\left[1 + \frac{3\sqrt{\pi}}{64}G_1 + \frac{3\sqrt{\pi}}{8}G_2\frac{\Delta d}{d_n}\right.$$
$$\left.+\left(2+\frac{51\sqrt{\pi}}{32}G_1\right)\Delta\alpha + (1-\alpha_n)\frac{3\sqrt{\pi}}{16}G_1\right]. \quad (44)$$

In the case of an inhomogeneous absorption line broadening in the free-molecular regime ($G_1 = G_2 = 0$), Eqs. (43) and (44) are simplified and, with allowance for Eq. (40), take the form

$$\frac{dp}{dz} = \frac{2\pi^{3/2}}{\lambda}\frac{\Omega}{k\bar{v}}\left(\frac{|G_{mn}|}{k\bar{v}}\right)^2 p_0\left\{\left[5 + 2\left(\frac{\Omega}{k\bar{v}}\right)^2\right]\Delta\alpha\right.$$
$$\left.+\frac{\hbar\omega_{mn}}{k_B T_0}(1+2\Delta\alpha)\right\}\exp\left[-\left(\frac{\Omega}{k\bar{v}}\right)^2\right], \quad (45)$$

$$\frac{dT}{dz} = \frac{4\pi^{3/2}}{\lambda}\frac{\Omega}{k\bar{v}}\left(\frac{|G_{mn}|}{k\bar{v}}\right)^2$$
$$\times T_0\left\{\left[-3 + 2\left(\frac{\Omega}{k\bar{v}}\right)^2\right]\Delta\alpha\right.$$
$$\left.+\frac{\hbar\omega_{mn}}{k_B T_0}(1+2\Delta\alpha)\right\}\exp\left[-\left(\frac{\Omega}{k\bar{v}}\right)^2\right]. \quad (46)$$

7. DISCUSSION OF RESULTS

Kinetic coefficients Λ_{pp}, Λ_{pT}, Λ_{TT}, and Λ_{Tp} determine the Poiseuille flow, thermal creep, heat flux, and mechanocaloric flow, respectively. In the Knudsen regime ($\delta \longrightarrow 0$), Eqs. (33)–(35) for these coefficients coincide fully with the well-known coefficients (see, e.g., [16]). In the collision regime, however, a dependence on radiative decay frequency Γ_m appears. This is related to the model of the gas, which represents a mixture of excited and unexcited particles, and particle exchange proceeds continuously between the gas components due to stimulated transitions, collisional transitions, and the radiative decay of the excited level.

Kinetic coefficients Λ_{pS}, Λ_{Sp}, Λ_{TS}, and Λ_{ST}, which are determined by Eqs. (36) and (37) and characterize the cross light-induced flows, represent odd functions of detuning Ω. Therefore, in an exact resonance at $\Omega = 0$, these coefficients are zero and the gas, heat, and entropy flows determined by them change their directions into opposite ones as the sign of detuning changes. Light-induced drift Λ_{pS} and cross baroentropy flow Λ_{Sp} are caused by the differences in the accommodation coefficients and the collision cross sections of excited and unexcited particles. Apart from these factors, light-induced heat transfer Λ_{TS} also includes excitation energy flux $\hbar\omega_{mn}$ [18].

Coefficient Λ_{SS}, which characterizes forward entropy flow $J_{SS} = \Lambda_{SS}X_S$, is independent of detuning in the case of the Doppler limit ($\Gamma \ll k\bar{v}$) in the zeroth approximation in parameter $\Gamma/(k\bar{v})$, is maximal for a homogeneous broadening ($\Gamma \gg k\bar{v}$) for exact resonance $\Omega = 0$, and decreases as the detuning increases.

The pressure and temperature gradients that form in the closed heat-insulated gas–capillary system are also odd functions of detuning. This means that the directions of the pressure and temperature gradients are inverted when detuning changes its sign.

We now estimate the pressure and temperature gradients that form in the free-molecular regime for CH_3Br vapors (transition wavelength $\lambda_{mn} \approx 10.6$ μm). Let the capillary radius be $r_0 = 1$ mm, the average system temperature be $T_0 = 500$ K, and the pressure be $p_0 = 0.1$ Pa. Under these conditions, the free-molecular regime takes place ($\delta = 0.01$) and the Doppler broadening is $k\bar{v} = 1.8 \times 10^8$ s^{-1}. We also assume that $|\Delta\alpha| = 0.005$, the Rabe frequency is $G_{mn} = 4 \times 10^4$ s^{-1}, and the detuning parameter is $\Omega/k\bar{v} = 0.3$. In this case, all assumptions related to the linearization of the problem are fulfilled.

From Eqs. (45) and (46), we obtain

$$\frac{dp}{dz} \approx 4 \times 10^{-3}\ \frac{\text{Pa}}{\text{m}}, \quad \left|\frac{dT}{dz}\right| \approx 40\ \frac{\text{K}}{\text{m}}.$$

The estimation of the pressure gradient shows that, in the free-molecular regime, it can be detected at the sensitivity limit of the existing pressure transducers. This finding is experimentally supported by the data obtained for CH_3F molecules [4]. The temperature gradient can be measured by any of the existing methods. As the Kn number decreases, collisions between gas particles are expected to equalize the gas temperature, i.e., to decrease the temperature gradient. Indeed, the excitation energy in the Knudsen gas transforms into heat at the end of the capillary, and this process occurs throughout the gas volume during collisions between gas particles.

Note that the contribution of the excitation energy flux to LIHT and the temperature gradient is approximately two orders of magnitude larger than the contribution of the flux of the translational energy of particles, which is proportional to $\Delta\alpha$. This behavior is likely to be typical of moderate temperatures in the optical, near, and medium IR wavelength ranges, at least for the Knudsen gas. It is seen from Eqs. (45) and (46) that the contribution of the excitation energy flux decreases in the far IR wavelength range at high temperatures, where $\hbar\omega_{mn}/k_BT \ll 1$, and it can even become smaller than the contribution of the translational energy flux under certain conditions. In this case, the pressure and temperature gradients are directed toward different sides in the almost free-molecular regime. The gas is cooled irrespective of the sign of detuning near the end of the tube, where the pressure is maximal. In contrast, the gas is heated near the other end, where the pressure is minimal.

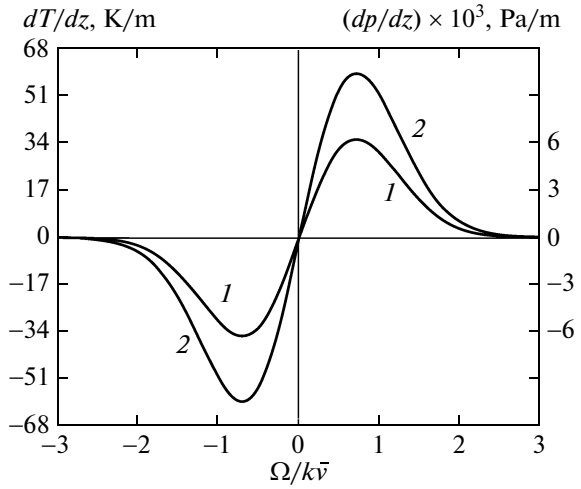

Fig. 3. (*1*) Pressure gradient and (*2*) temperature gradient vs. detuning parameter $\Omega/k\bar{v}$.

Figure 3 shows the pressure gradient (curve *1*) and the temperature gradient (curve *2*) versus detuning parameter $\Omega/k\bar{v}$ that were calculated for the free-molecular regime by Eqs. (45) and (46), respectively. The pressure and temperature gradients are seen to reach their maximum absolute values at $|\Omega/k\bar{v}| \approx 0.7$ and almost vanish at $|\Omega/k\bar{v}| \geq 2.5$, where the resonance condition is violated.

In conclusion, note that the theory and numerical estimates were made in the approximation of an infinitely long capillary, where the effect of the heat conduction of its wall may be neglected. In the case of a capillary of a finite length, the temperature distribution over its surface should be calculated. To this end, it is necessary to solve a thermophysical problem, as was proposed in [11] and was done in [18] under certain approximations. In this case, the temperature drop in the gas at the ends depends on the thermal conductivity of the capillary material. Moreover, the flows and the excitation energy dissipation can be cor-

rectly described near the ends of a capillary of a finite length.

REFERENCES

1. F. Kh. Gel'mukhanov and A. M. Shalagin, Sov. Phys. JETP **51** (5), 839 (1980).
2. A. V. Gainer, K. P. Komarov, and K. G. Folin, Sov. Phys. JETP **55** (6), 1068 (1982).
3. I. V. Chermyaninov, V. G. Chernyak, and E. A. Vilisova, JETP **105** (3), 511 (2007).
4. R. W. M. Hoogeveen, G. J. Van der Meer, and L. J. F. Hermans, Phys. Rev. A: At., Mol., Opt. Phys. **42**, 6471 (1990).
5. S. R. De Groot and P. Mazur, *Non-Equilibrium Thermodynamics* (North-Holland, Amsterdam, The Netherlands, 1962; Mir, Moscow, 1964).
6. F. Sharipov, J. Stat. Phys. **78**, 413 (1995).
7. V. G. Chernyak and A. P. Polikarpov, J. Stat. Phys. **140**, 504 (2010).
8. F. Kh. Gel'mukhanov and L. V. Il'ichev, Khim. Fiz., No. 11, 1544 (1984).
9. C. Cercignani, *Theory and Application of the Boltzmann Equation* (Plenum, New York, 1975; Mir, Moscow, 1978).
10. L. D. Landau and E. M. Lifshitz, *Course of Theoretical Physics*, Volume 10: E. M. Lifshitz and L. P. Pitaevskii, *Physical Kinetics* (Nauka, Moscow, 1979; Butterworth–Heinemann, Oxford, 1981).
11. A. M. Bishaev and V. A. Rykov, Izv. Akad. Nauk SSSR, Mekh. Zhidk. Gaza, No. 6, 106 (1984).
12. F. Sharipov, Physica A (Amsterdam) **203**, 457 (1994).
13. S. K. Loyalka, J. Chem. Phys. **71**, 339 (1979).
14. D. P. Willis, in *Rarefied Gas Dynamics*, Ed. by F. M. Devienne (Pergamon, New York, 1960; Inostrannaya Literatura, Moscow, 1963).
15. P. L. Chapovsky, G. J. Van der Meer, J. Smeets, and L. J. F. Hermans, Phys. Rev. A: At., Mol., Opt. Phys. **45**, 8011 (1992).
16. I. V. Chermyaninov, V. G. Chernyak, and G. A. Fomyagin, Teplofiz. Vys. Temp. **23**, 1158 (1985).
17. B. D. Fried and S. D. Conte, *The Plasma Dispersion Function* (Academic, New York, 1961).
18. P. L. Chapovskii and A. M. Shalagin, Kvantovaya Elektron. (Moscow) **12**, 227 (1985).

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